hep-ph/9604314
April 1996

Strange and Charmed Quarks in the Nucleon.

Ariel R. Zhitnitsky[1]

Physics Department, University of British Columbia, 6224 Agricultural Road, Vancouver, BC V6T 1Z1, Canada

Abstract

We discuss the general method of the calculation of the nucleon matrix elements of an operator associated with nonvalence quarks. The method is based on the QCD sum rules and low energy theorems. As an application of these considerations, we calculate the strange quark matrix element as well as the momentum distribution of the strangeness in the nucleon. We also calculate the singlet axial constant associated with η' meson as well as an axial constant associated with heavy quarks.

Talk given at the Workshop on the Structure of the η' Meson. NMSU and CEBAF, Las Cruces, New Mexico, March 8,1996.

[1] e-mail address:arz@physics.ubc.ca

1 Introduction

For a long time it was widely believed that the admixture of the pairs of strange quarks in the nucleons is small. The main justification of this picture was the constituent quark model where there is no room for strange quark in the nucleon. It has been known for a while that this picture is not quite true: In scalar and pseudoscalar channels one can expect a noticeable deviation from this naive prediction. This is because, these channels are very unique in a sense that they are tightly connected to the QCD-vacuum fluctuations with $0^+, 0^-$ singlet quantum numbers. Manifestation of the uniqueness can be seen, in particular, in the existence of the axial anomaly (0^- channel) and the trace anomaly (0^+ channel). Nontrivial QCD vacuum structure tells us that one could expect some unusual properties when we deal with those quantum numbers.

As we now know , this is indeed the case. In particular, we know that the strange quark matrix element $\langle N|\bar{s}s|N\rangle$ does not vanish and has the same order of magnitude as $\langle N|\bar{d}d|N\rangle$. This information can be obtained from the analysis of the so-called σ term [1],[2]. Similarly the analysis of the " proton spin crisis" essentially teaches us that the spin which is carried by the strange quark in the nucleon is not small as naively one could expect, see e.g. the recent review [3].

Another phenomenological manifestation of the same kind is the very old observation that in the scalar and pseudoscalar channels the Zweig rule is badly broken and there is substantial admixture of s quarks in the scalar mesons $f_0(980)$ (was S^*), $a_0(980)$ (was δ), $f_0(1300)$ (was ϵ), as well as in the pseudoscalar mesons η and η'. At the same time, in the vector channel the Zweig rule works well. Phenomenologically it is evident in e.g. the smallness of the $\phi - \omega$ mixing. In terms of QCD such a smallness corresponds to the numerical suppression of the nondiagonal correlation function $\int dx \langle 0|T\{\bar{s}\gamma_\mu s(x), \bar{u}\gamma_\nu u(0)\}|0\rangle$ in comparison with the diagonal one $\int dx \langle 0|T\{\bar{u}\gamma_\mu u(x), \bar{u}\gamma_\nu u(0)\}|0\rangle$. In the scalar and pseudoscalar channels diagonal and non-diagonal channels have the same order of magnitude. We believe that analysis of such kind of the correlation functions is an appropriate method for a QCD- based explanation of the unusual hadronic properties mentioned above.

In this talk we present some general methods and ideas for the analysis of the nucleon matrix elements from a non-valence operator. The ideology and methods (unitarity, dispersion relations, duality, low-energy theorems) we use are motivated by QCD sum rules. However we do not use the QCD sum rules in the common sense. Instead, we reduce one complicated problem (the calculation of non-valence nucleon matrix elements) to another one (the behavior of some vacuum correlation functions at low momentum transfer). One could think that such a reducing of one problem to another one (may be even more complicated) does not improve our understanding of the

phenomenon. However, this is not quite true: The analysis of the vacuum correlation functions with vacuum quantum numbers, certainly, is a very difficult problem. However some nonperturbative information based on the low energy theorems is available for such a correlation function. This gives some chance to estimate some interesting quantities.

2 Strangeness in the nucleon, 0^+ channel.

2.1 First estimations

We start by calculating the strange scalar matrix elements over the nucleon, assuming an octet nature of $SU(3)$ symmetry breaking. We follow to ref.[4](see also the book [5] for a review) in our calculations [6], but with a small difference in details. We present these results for completeness of the talk.

The results of the fit to the data on πN scattering presented in [2] lead to the following estimates for the so-called σ term

$$\frac{m_u+m_d}{2}\langle p|\bar{u}u+\bar{d}d|p\rangle=(64\pm 8MeV). \tag{1}$$

(Here and in what follows we omit kinematical structure like $\bar{p}p$ in expressions for matrix elements.) Taking the values of quark masses to be $m_u=5.1\pm 0.9MeV$, $m_d=9.3\pm 1.4MeV$, $m_s=175\pm 25MeV$ [7], from (1) we have

$$\langle p|\bar{u}u+\bar{d}d|p\rangle\simeq 9. \tag{2}$$

Further, assuming octet-type $SU(3)$ breaking to be responsible for the mass splitting in the baryon octet, we find

$$\langle p|\bar{u}u-\bar{d}d|p\rangle=\frac{m_\Xi-m_\Sigma}{m_s}=0.7, \tag{3}$$

$$\langle p|\bar{u}u+\bar{d}d-2\bar{s}s|p\rangle=3\frac{m_\Xi-m_\Lambda}{m_s}=3.4. \tag{4}$$

Here $m_\Xi, m_\Sigma, m_\Lambda$ are masses of Ξ, Σ, Λ hyperons respectively. The values (3), (4) are quite reasonable: the former is close to the difference of the number of u and d quarks in a proton (should be 1), and the latter is close to the total number of valence quarks u and d in a nucleon (see below). from (2-4) one obtains:

$$\langle p|\bar{u}u|p\rangle\simeq 4.8, \tag{5}$$

$$\langle p|\bar{d}d|p\rangle\simeq 4.1, \tag{6}$$

$$\langle p|\bar{s}s|p\rangle\simeq 2.8. \tag{7}$$

We should mention that the accuracy of these equations is not very high. For example, the error in the value of the σ term already leads to an error of order of one in each matrix element discussed above. However, these very simple calculations explicitly demonstrate that the strange matrix element is by no means small.

We would like to rewrite the relations (5-7) to separate the vacuum contribution to the nucleon matrix element from the valence contribution. In order to do so, let us define

$$\langle p|\bar{q}q|p\rangle \equiv \langle p|\bar{q}q|p\rangle_0 + \langle p|\bar{q}q|p\rangle_1, \tag{8}$$

where index 0 labels a (sea) vacuum contribution and index 1 a valence contribution for a quark q. We assume that the vacuum contribution which is related to the sea quarks is the same for all light quarks u, d, s. Thus, the nonzero magnitude for the strange matrix elements comes exclusively from the vacuum fluctuations. At the same time, the matrix elements related to the valence contributions are equal to

$$\langle p|\bar{u}u|p\rangle_1 \simeq (4.8 - 2.8) \simeq 2, \tag{9}$$

$$\langle p|\bar{d}d|p\rangle_1 \simeq (4.1 - 2.8) \simeq 1.3. \tag{10}$$

These values are in remarkable agreement with the numbers 2 and 1, which one could expect from the naive picture of non-relativistic constituent quark model. In spite of the very rough estimations presented above, we believe we convinced a reader that :

a) a magnitude of the nucleon matrix element for $\bar{s}s$ is not small;

b) the large magnitude for this matrix element is due to the nontrivial QCD vacuum structure where vacuum expectation values of u, d, s quarks are developed and they are almost the same in magnitude: $\langle 0|\bar{d}d|0\rangle \simeq \langle 0|\bar{u}u|0\rangle \simeq \langle 0|\bar{s}s|0\rangle$.

Once we realized that the phenomenon under discussion is related to the nontrivial vacuum structure, it is clear that the best way to understand such a phenomenon is to use some method where QCD vacuum fluctuations and hadronic properties are strongly interrelated. We believe, that the most powerful analytical nonperturbative method which exhibits these features is the QCD sum rules approach [8],[9].

In what follows we use the QCD sum rule method in order to relate hadronic matrix elements and vacuum characteristics. Let me emphasize from the very beginning that we do not use the QCD sum rules in the standard way: we do not fit them to extract any information about lowest resonance (as people usually do in this approach), we do not use any numerical approximation or implicit assumptions about higher states. Instead, we concentrate on the qualitative relations between hadronic properties and QCD vacuum structure. We try to explain in qualitative way some magnitudes

for the nucleon matrix elements which may look very unexpected from the naive point of view. At the same time those matrix elements can be easily understood in terms of the QCD vacuum structure.

We close this section with the formulation of the following question:
Q: What is the QCD explanation of the unusual properties mentioned above? (in particular, the large magnitude for the strange nucleon matrix element, a special role of the scalar and pseudoscalar channels et cetera). Our answer on this question is:
A: Hadronic matrix elements with $0^{\pm}$ *quantum numbers are singled out because of the special role they play in the QCD vacuum structure.* The next section switches this answer from a qualitative remark into the quantitative description.

2.2 Strangeness in the nucleon and vacuum structure

To study the problem of calculation $\langle N|\bar{s}s|N\rangle$ using the QCD -sum rules approach, we consider the following vacuum correlation function [6]:

$$T(q^2) = \int e^{iqx} dx dy \langle 0|T\{\eta(x), \bar{s}s(y), \bar{\eta(0)}\}|0\rangle \qquad (11)$$

at $-q^2 \to \infty$. Here η is an arbitrary current with nucleon quantum numbers. In particular, this current may be chosen in the standard form $\eta = \epsilon^{abc}\gamma_\mu d^a(u^b C\gamma_\mu u^c)$. Note, however, that the results obtained below do not imply such a concretization. For the future convenience we consider the unit matrix kinematical structure in (11).

This is the standard first step of any calculation of such a kind: Instead of direct calculation of a matrix element, we reduce the problem to the computation of some correlation function. As the next step, we use the duality and dispersion relations to relate a physical matrix element to the QCD-based formula for the corresponding correlation function. This is essentially the basic idea of the QCD sum rules.

In our specific case (11) due to the absence of the s -quark field in the nucleon current η, any substantial contribution to $T(q^2)$ is connected only with non-perturbative, so-called induced vacuum condensates, see Fig.1. Such a contribution arises from the region, when some distances are large: $(y-0)^2 \sim (y-x)^2 \gg (x-0)^2$. Thus, it can not be directly calculated in perturbative theory, instead we code the corresponding large-distance information in the form of a bilocal operator

$$K = i \int dy \langle 0|T\{\bar{s}s(y), \bar{u}u(0)\}|0\rangle, \qquad (12)$$

see Fig.1 The similar contributions were considered at first time in ref.[10]. For the different applications of this approach when bilocal operators play essential role, see also refs.[11].

Along with consideration of the three-point correlation function (11), we would like to consider the standard two-point correlator

$$P(q^2) = \int e^{iqx} dx \langle 0 | T\{\eta(x), \bar{\eta(0)}\} | 0 \rangle, \qquad (13)$$

see Fig.2. The correlator (13) is determined by the nucleon residues $\langle 0|\eta|N\rangle$ and some duality interval S_0. At the same time the correlator (11) includes the information on the nucleon matrix element $\langle N|\bar{s}s|N\rangle$ also. Comparing (11) with (13) at $-q^2 \to \infty$, we arrive to the following relation [6]:

$$\langle N|\bar{s}s|N\rangle \simeq \frac{-m}{\langle \bar{q}q\rangle} K, \qquad (14)$$

where m is the nucleon mass. The main assumptions which have been made in the derivation of this relation are the following. First, we made the standard assumption about local duality for the nucleon. In different words we assumed that a nucleon saturates both correlation functions with duality interval S_0. The second assumption is that the typical scales (or what is the same, duality intervals in the limit $-q^2 \to \infty$ in the corresponding sum rules ((11) and (13)) are not much different in magnitude from each other. In this case the dependence on residues $\langle 0|\eta|N\rangle$ is canceled out in the ration and we are left with the matrix element $\langle N|\bar{s}s|N\rangle$ (14) we are interested in.

Note, that both these assumptions are very likely to be satisfied because we know that in most cases the lowest state (nucleon) does saturate the sum rules. If it does, than the typical scale (which in variety of sum rules is one and the same and of order of $1GeV^2$) guarantees that the duality intervals are likely to be very close to each other. Anyway, the quantitative analysis of the corresponding sum rules is possible, however it is not our main goal; rather we want to demonstrate the relation between matrix elements like $\langle N|\bar{s}s|N\rangle$ and the corresponding vacuum properties which are hidden in the correlator K (12). In principle one could analyze the sensitivity of the corresponding QCD sum rules to the lowest state, nucleon. Once it is demonstrated, we believe that the accuracy of our formula (14) is of order $20\% - 30\%$ which is a typical error for the sum rule approach.

Thus, the calculating of $\langle N|\bar{s}s|N\rangle$ reduces to the evaluation of the vacuum correlator K. Fortunately, sufficient information about the latter comes from the low -energy theorems. We note also that this method of reducing the nucleon matrix elements to that of the vacuum correlator is directly generalized to cover the arbitrary scalar O_S or pseudoscalar O_P operator[2]:

$$\langle N|O_S|N\rangle \simeq \frac{-m\bar{N}N}{\langle \bar{q}q\rangle} i \int dy \langle 0|T\{O_S, \bar{u}u(0)\}|0\rangle, \qquad (15)$$

[2]We assume of course that these operators do not contain u, d quarks. Otherwise an additional contribution which comes from the small distances must be also included.

$$\langle N|O_P|N\rangle \simeq \frac{-m\bar{N}i\gamma_5 N}{\langle\bar{q}q\rangle} i\int dy\langle 0|T\{O_P, \bar{u}i\gamma_5 u(0)\}|0\rangle. \tag{16}$$

The estimation of the nonperturbative correlator K can be done by using some low-energy theorems. In this case K is expressed in terms of some vacuum condensates [6]:

$$K = i\int dy\langle 0|T\{\bar{s}s(y), \bar{u}u(0)\}|0\rangle \simeq \frac{18}{b}\frac{\langle\bar{q}q\rangle^2}{\langle\frac{\alpha_s}{\pi}G^2_{\mu\nu}\rangle} \simeq 0.04GeV^2, \tag{17}$$

where $b = \frac{11}{3}N_c - \frac{2}{3}N_f = 9$ and we use the standard values for the vacuum condensates[8]:

$$\langle\frac{\alpha_s}{\pi}G^2_{\mu\nu}\rangle \simeq 1.2\cdot 10^{-2}GeV^4 \quad \langle\bar{q}q\rangle \simeq -(250MeV)^3.$$

With the estimation (17) for K, our formula (14) gives the following expression for the nucleon expectation value for $\bar{s}s$

$$\langle p|\bar{s}s|p\rangle \simeq -m\cdot\frac{18}{b}\frac{\langle\bar{q}q\rangle}{\langle\frac{\alpha_s}{\pi}G^2_{\mu\nu}\rangle} \simeq 2.4, \tag{18}$$

which is very close to the naive estimation (7). Let us stress: we are not pretending to have made a reliable calculation of the matrix element $\langle p|\bar{s}s|p\rangle$ here. Rather, we wanted to emphasize on the qualitative picture which demonstrates the close relation between nonvalence matrix elements and QCD vacuum structure.

We close this section by noting that the method presented above gives very simple physical explanation of why the Zweig rule in the scalar and pseudoscalar channels is badly broken and at the same time, in the vector channel the Zweig rule works well. In particular, the matrix element $\langle N|\bar{s}\gamma_\mu s|N\rangle$ is expected to be very small as well as the corresponding coupling constant $g_{\phi NN}$ does. In terms of QCD such a smallness corresponds to the numerical suppression $(10^{-2} - 10^{-3})$ of the nondiagonal correlation function $\int dx\langle 0|T\{\bar{s}\gamma_\mu s(x), \bar{u}\gamma_\nu u(0)\}|0\rangle$ in comparison with the diagonal one $\int dx\langle 0|T\{\bar{u}\gamma_\mu u(x), \bar{u}\gamma_\nu u(0)\}|0\rangle$, see QCD-estimation in [8]. In the scalar and pseudoscalar channels the diagonal and non-diagonal channels have the same order of magnitude.

In the next few sections we discuss some applications of the obtained results.

2.3 In the world where s quark is massless.

We would like to look at formula (17) from the different side. Namely, we note that K not only enters expression (14), but also determines the variation of the condensate $\langle\bar{u}u\rangle$ with s quark mass:

$$\frac{d}{dm_s}\langle\bar{u}u\rangle = -i\int dy\langle 0|T\{\bar{s}s(y), \bar{u}u\}|0\rangle = -K \simeq -0.04GeV^2. \tag{19}$$

To understand how large this number is and in order to make some rough estimations, we assume that this behavior can be extrapolated from physical value $m_s \simeq 175MeV$ till $m_s = 0$. In this case we estimate that

$$| \frac{\langle \bar{u}u \rangle_{m_s=175} - \langle \bar{u}u \rangle_{m_s=0}}{\langle \bar{u}u \rangle_{m_s=175}} | \simeq 0.5. \tag{20}$$

Such a decrease of $| \langle \bar{u}u \rangle |$ by a factor of two as m_s varies from $m_s \simeq 175MeV$ to $m_s = 0$ is a very important consequence of the previous discussions: Once we accept the relatively large magnitude for the nucleon matrix element $\langle p|\bar{s}s|p\rangle \simeq 2.4$, we are forced to accept the relatively large variation of the light quark condensate as well. This statement is the direct consequence of QCD, see (20).

We note that this result does not seem very surprising since other vacuum condensates, e.g. $\langle \frac{\alpha_s}{\pi} G^2_{\mu\nu} \rangle$ possess analogous properties [12]. From the microscopic point of view, decrease of absolute values of vacuum matrix elements with the decrease of the s quark mass is expected since any topologically nontrivial vacuum configurations, e.g. instantons, are suppressed by light quarks. The corresponding numerical calculation is very difficult to perform, however a qualitative picture of the QCD vacuum structure definitely supports this idea [13].

2.4 s quark and the nucleon mass.

We would like to discuss here one more fundamental characteristic of the hadron world: the nucleon mass and its dependence on the strange quark. We start our discussion from the following well known result: the nucleon mass is determined by the trace of the energy -momentum tensor $\theta_{\mu\mu}$ and in the chiral limit $m_u = m_d = m_s = 0$ the nonzero result comes exclusively from the strong interacting gluon fields:

$$m = -\frac{b}{8} \langle N | \frac{\alpha_s}{\pi} G^2_{\mu\nu} | N \rangle, \quad m_u = m_d = m_s = 0. \tag{21}$$

However, as we know, in our world the strange quark is not massless, but rather it requires some (large enough) mass ($\sim 175MeV$). As we have seen (20), the nonzero mass of s quark considerably changes the vacuum properties of the world. Thus, we would expect that it might have strong influence on the nucleon mass as well. The main argument which supports this point of view is the same as before, and is based on our general philosophy that the nucleon matrix elements and vacuum properties are tightly related. So, if the strange quark has strong influence on the vacuum properties than its impact on the nucleon mass should also be strong.

In order to check these reasons it would be useful to calculate the strange quark contribution into the nucleon mass directly and independently from the gluon contribution (21). Fortunately, it can be easily done by using our

previous estimation (18) for the nucleon matrix element and exact expression for the trace of the energy-momentum tenzor with taking into account nonzero quark masses:

$$m = +\langle N|\sum_q m_q\bar{q}q|N\rangle - \frac{b}{8}\langle N|\frac{\alpha_s}{\pi}G^2_{\mu\nu}|N\rangle, \qquad (22)$$

where sum over q is sum over all light qurks u, d, s. One can easily see from (1) that u, d contribution into the nucleon mass does not exceed 7%; thus we can safely neglect this. At the same time, adopting the values (7),(18) for $\langle p|\bar{s}s|p\rangle$ and $m_s \simeq 175MeV$ [7], one can conclude that considerable part of the nucleon mass (about 45%) is due to the strange quark. In this case the gluon contribution into the nucleon mass is far away from the chiral $SU(3)$ prediction (21) and approximately equals to

$$-\frac{b}{8}\langle N|\frac{\alpha_s}{\pi}G^2_{\mu\nu}|N\rangle \sim 520MeV. \qquad (23)$$

This rough estimation confirms our argumentation that a variation of the strange quark mass from its physical value to zero, may considerably change some vacuum characteristics as well as nucleon matrix elements.

The simple consequence of this result is the observation that the **quenched approximation** in the lattice calculations is not justified simply because such a calculation clearly not accounting the fluctuations of the strange (non-valence) quark as well as vacuum fluctuations of u and d quarks. As we argued above, the nucleon mass undergoes some influence from s quark.

How one can understand these results within the framework of the QCD sum rules? Let us recall that in the QCD sum rules approach an information about any dimensional parameter is contained in the vacuum condensates $\langle\bar{u}u\rangle, \langle G^2_{\mu\nu}\rangle, \ldots$. As we discussed previously all these condensates varying with m_s considerably. It is important that this variation certainly proceeds in the right direction: Absolute values of condensates decrease with decreasing m_s. This leads to a smaller scale in the sum rules and finally, to the decrease of all dimensional parameters such as m. However, it is difficult to make any reliable calculations because of a large number of factors playing an essential role in such a calculation.

2.5 Momentum distribution of the strangeness in the nucleon .

We continue our study on the role of the strange quark in nucleon with the following remark. We found out earlier that the matrix element $\langle N|\bar{s}s|N\rangle$ is not small; we interpreted this result as a result of strong vacuum fluctuations which penetrate into the nucleon matrix element. Now, we would like to ask the following question: What is the mean value of the momentum (denoted

as $\langle k_\perp^2\rangle_s$) of the s quark inside of a nucleon? Let us note that this question is not a pure academic one. Rather, the answer on the question might be important for the construction of a more sophisticated quark model which would incorporate the strange context into the nucleon wave function.

First of all, let us try to formulate this question in terms of QCD. We *define* the mean value $\langle k_\perp^2\rangle_s$ of the momentum carried by the strange quark in a nucleon by the following matrix element:

$$\langle k_\perp^2\rangle_s\langle N|\bar{s}s|N\rangle \equiv \langle|N\bar{s}(i\overrightarrow{D}_\perp)^2 s|N\rangle, \tag{24}$$

where $i\overrightarrow{D}_\mu \equiv i\overrightarrow{\partial}_\mu + gA_\mu^a\frac{\lambda^a}{2}$ is the covariant derivative and A_μ^a is gluon field. The arrow shows the quark whose momentum is under discussion.

We assume that the nucleon to be moving rapidly in the z-direction. We are interested in the momentum distribution in the direction which is perpendicular to its motion. Precisely this characteristic has a dynamical origin. Indeed, as we shall see in a moment, while we are studying a nucleon matrix element $\langle k_\perp^2\rangle_s$, we are actually probing the QCD vacuum properties! The nucleon motion as a whole system with arbitrary velocity does not affect this characteristic. Thus, essentially, what we discuss is the, so-called, light cone wave function. Apart of the reasons mentioned above there are few more motives to do so: First of all, the light cone wave function (wf) with a minimal number of constituents is a good starting point. As is known such a function gives the parametrically leading contributions to hard exclusive processes. Higher Fock states are also well defined in this approach and can be considered separately. The second reason to work with a light cone wave function is there existence of the nice relation between that wf and structure function measured in the deep-inelastic scattering. We refer to the review paper [14] for the introduction into the subject. The relation to the standard quark model wave functions (see e.g.[15]) is also worked out. The relevant discussions can be found in ref.[16]. Besides of these, we have one more reason to work with the light cone wf: we believe that this is the direction where a valence quark model can be understood and formulated in the QCD -terms[17].

Anyhow, the formula (24) with the derivatives taken in the direction perpendicular to the nucleon momentum $p_\mu = (E, 0_\perp, p_z)$, is very natural definition for the mean square of the quark transverse momentum. Of course it is different from the naive, gauge dependent definition like $\langle N|\bar{s}\partial_\perp^2 s|N\rangle$, because the physical transverse gluon is participant of this definition. However, the expression (24) is the only possible way to define the $\langle k_\perp^2\rangle_s$ in the gauge theory like QCD. We believe that such definition is the useful generalization of the transverse momentum conception for the interactive quark system. Let us note that the Lorentz transformation in z direction does not affect the transverse directions. Thus, the transverse momentum $\langle k_\perp^2\rangle_s$ calculated from eq. (24) remains unchanged while we passing from the light

cone system to the rest frame system where a quark model suppose to be formulated.

Now, let us come back to our definition (24) for $\langle k_\perp^2\rangle_s$. In order to calculate this matrix element, we use the same trick as before: we reduce our original problem of the calculating of a nucleon matrix element to the problem of a computing of the corresponding vacuum correlation function (15):

$$\langle N|\bar{s}(i\overrightarrow{D}_\perp)^2 s|N\rangle \simeq \frac{-m\bar{N}N}{\langle\bar{q}q\rangle} i\int dy\langle 0|T\{\bar{s}(i\overrightarrow{D}_\perp)^2 s, \bar{u}u(0)\}|0\rangle. \qquad (25)$$

To estimate the right hand side of the eq.(25) we introduce an auxiliary vacuum correlation function

$$i\int dy\langle 0|T\{\bar{s}(i\overrightarrow{D}_\mu\, i\overrightarrow{D}_\nu)s, \bar{u}u(0)\}|0\rangle = Cg_{\mu\nu}, \qquad (26)$$

where C is constant. From the definition (25) it is clear that the correlator we are interested in can be expressed in terms of the constant C:

$$i\int dy\langle 0|T\{\bar{s}(i\overrightarrow{D}_\perp)^2 s, \bar{u}u(0)\}|0\rangle = -2C. \qquad (27)$$

At the same time the constant C is given by the correlation function which contains $G^a_{\mu\nu}$ and not a covariant derivative D_μ:

$$C = \frac{1}{4} i\int dy\langle 0|T\{\bar{s}(i\overrightarrow{D}_\mu\, i\overrightarrow{D}_\mu)s, \bar{u}u(0)\}|0\rangle =$$

$$\frac{-1}{8} i\int dy\langle 0|T\{\bar{s}igG^a_{\mu\nu}\frac{\lambda^a}{2}\sigma_{\mu\nu}s, \bar{u}u(0)\}|0\rangle, \qquad (28)$$

where we have used the equation of motion and identity[3]:

$$D_\mu D_\nu g_{\mu\nu}s = \gamma_\mu\gamma_\nu D_\mu D_\nu s - \sigma_{\mu\nu}\frac{1}{2}[D_\mu, D_\nu]s = -m_s^2 s + \frac{ig}{2}\sigma_{\mu\nu}G^a_{\mu\nu}\frac{\lambda^a}{2}s \qquad (29)$$

Now we can estimate the unknown vacuum correlator (28) exactly in the same way as we have done before for the correlation function K, see eq.(17). Collecting all formulae (24-29) together, we arrive to the following final result for the mean value of the momentum carried by the strange quark in a nucleon:

$$\langle k_\perp^2\rangle_s \equiv \frac{\langle N|\bar{s}(i\overrightarrow{D}_\perp)^2 s|N\rangle}{\langle N|\bar{s}s|N\rangle} \simeq \frac{\langle N|\bar{s}ig\sigma_{\mu\nu}G^a_{\mu\nu}\frac{\lambda^a}{2}s|N\rangle}{4\langle N|\bar{s}s|N\rangle}$$

$$\simeq \frac{\langle\bar{s}ig\sigma_{\mu\nu}G^a_{\mu\nu}\frac{\lambda^a}{2}s\rangle}{4\langle\bar{s}s\rangle}\cdot\frac{d^{\bar{s}ig\sigma_{\mu\nu}G^a_{\mu\nu}\frac{\lambda^a}{2}s}}{d^{\bar{s}s}} \simeq \frac{1}{4}(0.8GeV^2)\frac{5}{3} \sim 0.33GeV^2, \qquad (30)$$

[3]We neglect the term proportional to m_s^2 in the eq.(28). It can be justified by using the estimation (17).

where d^O denotes the dimension of the operator O. For numerical estimation we use the standard magnitude for the mixed vacuum condensate $\langle\bar{s}ig\sigma_{\mu\nu}G^a_{\mu\nu}\frac{\lambda^a}{2}s\rangle = 0.8GeV^2\langle\bar{s}s\rangle$. The obtained numerical value (30) for $\langle k_\perp^2\rangle_s$ looks very reasonable from the phenomenological point of view.

We close this section with a few remarks. First, the nonvalence nucleon matrix elements can be expressed in terms of vacuum condensates in a very nice way. All numerical results obtained in such a way look very reasonable. As the second remark, we emphasize that a study of nonvalence nucleon matrix elements and an analysis of the QCD vacuum structure is one and the same problem. We would like to note also, that the nucleon matrix element (30) might be very important in the analysis of neutron dipole moment. This observation is based on the fact that the so called chromoelectric dipole moment of the s quark, related to the operator $\bar{s}g\gamma_5\sigma_{\mu\nu}G^a_{\mu\nu}\frac{\lambda^a}{2}s$, in many models gets a large factor $\sim m_s/m_q \sim 20$ in comparison with a similar d quark contribution [18], [6], [19]. At the same time, as we can see from (30) there is no any suppression due to the presence of the s quark in the corresponding nucleon matrix elements.

3 Strangeness in the nucleon, 0^-channel.

3.1 Singlet axial constant g^0_A.

In this section we discuss the contribution of the strange quarks into the nucleon matrix elements similar to eq. (22), with the only difference that we switch the scalar channel $\bar{s}s$ into the pseudoscalar one $\bar{s}i\gamma_5 s$. In our previous study of the scalar channel we concluded that the considerable part of the nucleon mass (about 40%) is due to the strange quark[4]. We made this estimation by using two following facts: First, we knew the mass of the nucleon (left hand side of the eq. (22)), which is considered as an experimental data. The second, we calculated independently the matrix element $\langle N|\bar{s}s|N\rangle$. Comparing this theoretical result (18) with (22), we have made aforementioned conclusion about a serious deviation from the chiral $SU(3)$ limit.

We want to repeat all these steps for the pseudoscalar channel also. In this case the equation analogous to (22) looks as follows:

$$2mg^0_A\bar{p}i\gamma_5 p = +\langle N|\sum_q 2m_q\bar{q}i\gamma_5 q|N\rangle + \frac{3}{4}\langle N|\frac{\alpha_s}{\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|N\rangle, \qquad (31)$$

where sum over q is the sum over all light qurks u, d, s and g^0_A is the nucleon axial constant in the flavor singlet channel. The world average is: $g^0_A = 0.27 \pm 0.04$, [3].

[4]In the chiral limit $m_s \to 0$, the corresponding contribution is zero, of course.

Now, we would like to repeat all steps which would bring us to the conclusion similar to eq.(22) for the pseudoscalar case. We shall try to answer on the following question: what is the strange quark contribution in the formula (31)? Let us recall that in the chiral limit $m_u = m_d = m_s = 0$ the nonzero contribution comes exclusively from the gluon term in the close analogy with formula (21):

$$2mg_A^0\bar{p}i\gamma_5 p = \frac{3}{4}\langle N|\frac{\alpha_s}{\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|N\rangle, \qquad m_u = m_d = m_s = 0. \tag{32}$$

Thus, in order to answer on the question formulated above, we have to estimate the matrix element

$$\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle \tag{33}$$

in somewhat independent way [5]. First of all, the relevant contribution with octet quantum numbers (η) can be easily evaluated by the standard technics. One should take the derivative from the octet, anomaly- free, current $\sim \bar{u}\gamma_\mu\gamma_5 u + \bar{d}\gamma_\mu\gamma_5 d - 2\bar{s}\gamma_\mu\gamma_5 s$. The result is:

$$\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle_\eta = -m(3F-D)\bar{p}i\gamma_5 p, \tag{34}$$

where $D \simeq 0.63g_A$ and $F \simeq 0.37g_A$ are the standard $SU(3)$ parameters. One could expect that the similar contribution with singlet quantum numbers (η') is also large, although it is zero in the chiral limit where $m_s = 0$.

We shall estimate the corresponding contribution with η'- quantum numbers by using our previous trick (16). Namely, we reduce our original problem of calculation of the nucleon matrix element to the problem of the computation of certain vacuum correlation function where we should limit ourself by calculating the contribution with singlet quantum numbers only:

$$\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle \simeq \frac{-m\bar{p}i\gamma_5 p}{\langle\bar{q}q\rangle} i\int dy\langle 0|T\{2m_s\bar{s}i\gamma_5 s, \bar{u}i\gamma_5 u(0)\}|0\rangle. \tag{35}$$

In order to make the corresponding estimations we need to know the following η'- matrix elements: $\langle 0|\bar{s}i\gamma_5 s|\eta'\rangle$ and $\langle 0|\bar{u}i\gamma_5 u|\eta'\rangle$. The PCAC does not provide us with the corresponding information, however a quark model prejudice suggests that

$$\langle 0|\frac{1}{\sqrt{2}}(\bar{u}i\gamma_5 u - \bar{d}i\gamma_5 d)|\pi\rangle \simeq \langle 0|\frac{1}{\sqrt{6}}(\bar{u}i\gamma_5 u + \bar{d}i\gamma_5 d - 2\bar{s}i\gamma_5 s)|\eta\rangle \simeq$$

$$\langle 0|\bar{u}i\gamma_5 s|K\rangle \simeq \langle 0|\frac{1}{\sqrt{3}}(\bar{u}i\gamma_5 u + \bar{d}i\gamma_5 d + \bar{s}i\gamma_5 s)|\eta'\rangle \simeq -\frac{2\langle\bar{q}q\rangle}{f_\pi} \tag{36}$$

[5] We neglect u, d qurk contributions into the formula (31) by the obvious reasons.

The strong support in favor that the relations (36) are to be correct, comes from the analysis of the two photon decays of π, η, η', see e.g.[5]. All of these decay amplitudes have the same Lorentz structure and determined by the matrix elements (36), therefore, the quark model prediction is found to work surprisingly well in this particular case. Combine the formulae (34-36) we arrive to the following estimation:

$$\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle = 2m\bar{p}i\gamma_5 p\{-\frac{1}{2}(3F-D) - \frac{4m_s\langle\bar{q}q\rangle}{3f_\pi^2 m_{\eta'}^2}\}$$

$$\simeq (-0.3+0.16)2m\bar{p}i\gamma_5 p \simeq (-0.14)2m\bar{p}i\gamma_5 p, \quad (37)$$

where we used $3F-D \simeq 0.6$ for the numerical estimation. The two terms in this formula are the octet and singlet contributions correspondingly. One should note, that in spite of the fact that the singlet term is parametrically suppressed in the limit $m_s = 0$, this contribution numerically is not small. It is only by a factor of two less than the parametrically leading term.

Now, let us come back to eq.(31). We would like to answer on the previously formulated question: what is the s quark contribution into the formula (31)? From our estimation (37) we suggest the following pattern of saturation of the experimental data for $g_A^0 = 0.27 \pm 0.04$:

$$2m\bar{p}i\gamma_5 p(0.27\pm 0.04) = +\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle + \frac{3}{4}\langle p|\frac{\alpha_s}{\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|p\rangle$$

$$\simeq (-0.14)2m\bar{p}i\gamma_5 p + (+0.41)2m\bar{p}i\gamma_5 p. \quad (38)$$

Here the first term is due to the strange quark contribution from (37) and the second one is due to the gluon contribution. We assign an average number 0.41 for the gluon contribution in order to match the experimental data for g_A^0.

Two remarks are in order. First, The strange-quark and the gluon terms contribute with the opposite signs into g_A^0. In the formula for mass (22) the similar terms interfere constructively, with the same signs. It is very easy to understand the difference: in the pseudoscalar channel we have the Goldstone boson, η, whose total contribution is zero into the sum (31) because of the octet origin of the η meson. However, the η meson contributions into the matrix elements $\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle$ and into the gluon operator, taken separately, are not zero. Moreover, its contribution to the $\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle$ has the opposite sign to the η' contribution (because of the difference in the quark context, see (36)). Even more, it has a parametrical enhancement. We have nothing like that in the scalar channel (22), where the flavor -singlet states dominate.

The second remark is the observation that, like in the scalar channel, the strange quark operator gives a noticeable contribution into the final formula (38) in spite of the fact that in the chiral limit the corresponding contribution is zero as we mentioned earlier (32).

3.2 Singlet axial constant of heavy particles.

In this section we try to answer on the following question: How one can check the estimation (38) about noticeable contribution of the strange quark? More specific question we would like to know: Is it possible to measure a nucleon matrix element where some independent combination of those operators enters? If answer were "yes", we would be able to find the contribution of each term separately.

To answer on this question we suggest to consider the weak neutral current containing an isoscalar axial component associated with nonvalence quarks [20]:

$$\langle p|\bar{c}\gamma_\mu\gamma_5 c - \bar{s}\gamma_\mu\gamma_5 s + \bar{t}\gamma_\mu\gamma_5 t - \bar{b}\gamma_\mu\gamma_5 b|p\rangle \equiv g_A^{heavy}\bar{p}\gamma_\mu\gamma_5 p \tag{39}$$

Before to go into details, let us mention, that on the quantum level the current divergence of the massive quark field has the following form:

$$\partial_\mu\bar{Q}\gamma_\mu\gamma_5 Q = 2m_Q\bar{Q}i\gamma_5 Q + \frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}, \tag{40}$$

where the first term is the standard one and the second term is due to the anomaly. There are many ways to understand the origin of the anomaly; basically it arises from the necessity of the ultraviolet regularization of the theory. In the heavy quark mass limit, one can expand $2m_Q\bar{Q}i\gamma_5 Q$ term in the eq.(40) with the following result [8],[9]:

$$2m_Q\bar{Q}i\gamma_5 Q|_{m_Q\to\infty} = -\frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G_{\mu\nu}} + c\frac{G\tilde{G}G}{m_Q^2} + 0(\frac{1}{m_Q^4}) + ... \tag{41}$$

where all coefficients, in principle, can be calculated. We are interested, however, in the leading term $\sim G_{\mu\nu}\tilde{G_{\mu\nu}}$ only.

One can easily note that the leading term in the expansion (41) has the same structure as an anomaly term (40) and it goes with the opposite sign[6]. Thus, the terms, which do not depend on mass are canceled out and we left with a term $\sim \frac{G\tilde{G}G}{m_Q^2}$ which vanishes in the limit $m_Q \to \infty$. Such a vanishing of the heavy quark contribution into the nucleon matrix element is in a perfect agreement with a physical intuition that a nucleon does not contain any heavy quark fields, at least in the limit $m_Q \to \infty$.

The situation with strange s quark in the formula (39) is much more complicated. This quark is not heavy enough to apply the arguments given

[6] The opposite signs of those contributions can be easily understood in terms of the Pauli-Villars regulator fields with mass $M_{PV} \to \infty$. As is known these fields are introduced into the theory for the regularization purposes and they play crucial role in the calculation of the anomaly (40). Regulator contribution is obtained, by definition, by a replacement $m_Q \to M_{PV}$ in the corresponding formula. It goes, by definition, with relative sign minus. From such a calculation it is clear that the leading terms which do not depend on mass, are canceled out in a full agreement with an explicit formula (40,41).

above. Thus, we should keep all operators in the formula (40) for the current divergence in the original form:

$$\langle p|2m_s\bar{s}i\gamma_5 s+\frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|p\rangle=-g_A^{heavy}2m\bar{p}i\gamma_5 p, \tag{42}$$

where we have neglected the terms $\sim\frac{1}{m_Q^2}$ for c,b,t quarks in according with our previous discussion[7].

As we already mentioned, the measurement of the constant g_A^{heavy} (42) gives an independent information complementary to the singlet axial constant measurement g_A^0, (38). If we knew those constants with high enough precision, we would be able to find out both nucleon matrix elements: $\langle p|2m_s\bar{s}i\gamma_5 s|p\rangle$ and $\langle p|\frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|p\rangle$. At the moment the experimental errors for $g_A^{heavy}=0.15\pm0.09$, [21] are large: the result is only two standard deviations from zero.

The best we can do at the moment is to estimate g_A^{heavy} from our previous calculations (38). If we literally take the values -0.14 and 0.41 from the formula (38), we get the result for g_A^{heavy} which is compatible with zero:

$$-g_A^{heavy}2m\bar{p}i\gamma_5 p=\langle p|2m_s\bar{s}i\gamma_5 s+\frac{\alpha_s}{4\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|p\rangle\simeq$$

$$\simeq(-0.14+\frac{1}{3}0.41)2m\bar{p}i\gamma_5 p,\quad g_A^{heavy}\simeq 0. \tag{43}$$

From our point of view this is an interesting observation which essentially says that the strange quark operator **together with its anomalous part** gives nearly **vanishing** contribution into the nucleon matrix element. As we discussed earlier, this is certainly true for any heavy quark. What is surprised us, that estimation (43) apperently says that this is true even for s quark (which is by no means can be considered as a heavy quark). If we accept this point, we should interpret a nonzero magnitude of g_A^0 (38) as a contribution coming exclusively from the light u,d quarks and their anomalous parts. As we mentioned, the s quark term together with its anomalous part gives almost vanishing contribution into eq.(38). Such an interpretation is in a very good agreement with the valence quark model philosophy, where the s quark does not play any essential role.

Let us note that this interpretation is very different from the old simplest assumption on the spin of the strange quark in the nucleon, see e.g.[3],[5]. In our interpretation we understand the strange quark contribution as a joined contribution of s field as well as its regulator field (or what is the same, its anomalous contribution).

[7] The corresponding estimations even for the lightest heavy c-quark support this viewpoint, see the next section.

4 The charmed quark in the nucleon.

In this section we would like to extend our analysis for the c quark. The reason to do so is twofold: First, the c quark is a heavy enough to use the standard $1/m_c$ expansion similar to (41). Secondly, the charmed quark is light enough to get a reasonably large effect from this expansion.

We start from the pseudoscalar channel and keep only the first term in the heavy quark expansion (41):

$$\langle p|\bar{c}i\gamma_5 c|p\rangle \simeq -\langle p|\frac{\alpha_s}{8m_c\pi}G_{\mu\nu}\tilde{G_{\mu\nu}}|p\rangle \sim \frac{-0.41}{6m_c}2m\bar{p}i\gamma_5 p \sim -0.1\bar{p}i\gamma_5 p \quad (44)$$

where, for the numerical estimate, we use the value (38) for the gluon matrix element over nucleon, and $m_c \simeq 1.3GeV$ for the charmed quark mass[8]. This value should be compared with the similar matrix element of the strange quark over nucleon:

$$\langle p|\bar{s}i\gamma_5 s|p\rangle \sim \frac{-0.14}{2m_s}2m\bar{p}i\gamma_5 p \sim -0.8\bar{p}i\gamma_5 p, \quad (45)$$

where we use formula (37) for the numerical estimation of the matrix element $\langle p|\bar{s}i\gamma_5 s|p\rangle$. The ratio of these values are in remarkable agreement with the ratio of their mass: $\frac{m_c}{m_s} \sim \frac{1.3GeV}{0.175GeV} \sim 7.5$.

Our next example is the scalar matrix element. In this case one can use the heavy quark expansion similar to formula (41), but for the scalar channel[9]:

$$\langle p|\bar{c}c|p\rangle \simeq -\langle p|\frac{\alpha_s}{12m_c\pi}G^a_{\mu\nu}G^a_{\mu\nu}|p\rangle \sim \frac{2\cdot 520MeV}{27m_c} \sim 0.03. \quad (46)$$

For the numerical estimation in this formula we adopted the value (23) for the gluon matrix element over nucleon. The magnitude (46) for the charmed quark is approximately hundred times less than the corresponding matrix element for the strange quark (18):

$$\langle p|\bar{s}s|p\rangle \simeq 2.4. \quad (47)$$

This is in a big contrast with pseudoscalar channel, where the corresponding ratio was about a factor ten larger.

We conclude this section with few remarks. First of all, the matrix elements (44),(46) for the charmed quark are expressed in terms of the gluon operators. For the heavy quark this is exact consequence of the QCD. Corrections to these formulae can be easily estimated. One can show that they

[8] Let us note, that this value for mass corresponds to the high enough normalization point of order of m_c. In principle, one should renormalize this value to the low normalization point. We neglect this small logarithmic effect in this paper.

are small for the c quark. The problem of the evaluation of the gluon matrix elements over nucleon is the different problem. However, we believe that from the measurements of g_A^0 (31) and from πN scattering (1) we know those matrix elements with a reasonable accuracy. Thus, we expect the same accuracy for the matrix elements $\langle p|\bar{c}c|p\rangle$ and $\langle p|\bar{c}i\gamma_5 c|p\rangle$.

Our next remark is the observation that results (44), (46) essentially give a normalization for the intrinsic charm quark component in the proton. This is very important characteristic of the nucleon. It might play an essential role in the explanation of a discrepancy between charm hadroproduction and perturbative QCD calculations. We refer to the original paper [22] [9] on this subject where the hypothesis of intrinsic charm quarks in the proton was introduced. The experimental fit in the framework of this paper[22] suggests that the probability to have an intrinsic charm in the proton is about $\sim 0.3\%$[24] and $(0.86\pm0.60)\%$[25]. These numbers can not be related directly to the matrix element (46) we calculated. However, they give some general scale of this phenomenon. We hope that in future, some more sophisticated, QCD-based methods, will lead us to deeper understanding of the effects related to intrinsic charm component in the nucleon.

5 Conclusion

We believe that the main result of the present analysis is the observation that non-valence quarks play an important role in the physics of nucleon. However we should stress that such an interpretation does not contradict to the bag model [5], where the nucleon matrix element

$$\langle N|\bar{s}s|N\rangle \simeq -\langle\bar{s}s\rangle V \tag{48}$$

is related to some vacuum characteristics (like condensate or volume of the bag V) of the model. It is clear that the chiral vacuum condensate of the strange quark is large. So, there is no reason to expect that the corresponding nucleon matrix element is small. The same argument can be applied for the arbitrary mixed vacuum condensates also (they are presumably not zero[17]). This information can be translated, in according to (30), into the knowledge about transverse momentum distribution of the strange quark in a nucleon.

In our approach this relation *vacuum* $\Longleftrightarrow$ *nucleon* is clearly seen. Thus, by studying the vacuum properties of the QCD, we essentially study some interesting nucleon matrix elements which can be experimentally measured.

References

[1] J.P.Cheng, Phys.Rev.**D 13**,(1976),2161

[9] See also recent paper [23].

[2] R.Koch Z.Phys.C-Particle and Fields, **15** (1982), 161.

[3] J.Ellis and M.Karliner, hep-ph/9601280, January, 1996.

[4] J.F.Donoghue, Ch.R.Nappi, Phys.Lett. **B 168**,(1986),105.

[5] J.F.Donoghue, E.Golowich and B.R.Holstein, Dynamics of the Standard Model, Cambridge University Press, 1992.

[6] V.M.Khatsimovsky,I.B.Khriplovich, A.R.Zhitnitsky, Z.Phys.C-Particle and Fields, **36** (1987), 455.

[7] H.Leutwyler, hep-ph/9602255, February, 1996.

[8] M.A.Shifman, A.I.Vainshtein and V.I.Zakharov Nucl.Phys. **B147**, (1979) 385, 448, 519.

[9] M.A.Shifman, Vacuum Structure and QCD Sum Rules, North-Holland, 1992.

[10] I.Balitsky, Phys.Lett. **B 114**,(1982), 53.

[11] I.Balitsky and A.Yung, Phys.Lett. **B 129**,(1983), 328.
B.L.Ioffe and A.Smilga ,Nucl.Phys. **B 232**, (1984), 109.
V.A.Nesterenko and A.V. Radyushkin, JETP Lett. **39**,(1984), 707.

[12] V.A.Novikov et al, Nucl.Phys. **B 191**, (1981), 301.

[13] E.Shuryak, Nucl.Phys. **B 203**, (1982), 93, 116.

[14] S.J.Brodsky and G.P.Lepage,"Exclusive Processes in Quantum Chromodynamics" in Perturbative Quantum Chromodynamics edited by A.H.Mueller, World Scientific Publishing Co.,1989.

[15] N.Isgur, in The New Aspects of Subnuclear Physics, edited by A.Zichichi (Plenum, New York, 1980);
J.Rosner, in Techniques and Concepts of High Energy Physics, edited by T.Ferbel (Plenum, New York, 1981).

[16] S.J.Brodsky, T.Huang and P. Lepage, in Particle and Fields, edited by A.Z. Capri and A.N.Kamal (Plenum Publishing Corporation, New Yourk, 1983).

[17] A. Zhitnitsky, Phys.Lett.**B357**,(1995),211;
B.Chibisov and A.Zhitnitsky, Phys.Rev.**D52**,(1995),5273.

[18] X.-G.He, B.H.J.McKellar, S.Pakvasa, Phys.Lett.**B254**,(1991),231; Int.J.Mod.Phys. **A4**, (1989), 5011.

[19] I.B.Khriplovich, hep-ph/9604212.

[20] J.Collins, F.Wilczek, A.Zee, Phys.Rev.**D18**,(1978),242.

[21] Ahrens et al. Phys.Rev.**D35**,(1987),785.

[22] S.Brodsky and C. Peterson, Phys.Rev.**D23**,(1981),2745.

[23] G.Ingelman and M.Thunman hep-ph/9604289, April 1996.

[24] E.Hoffman and R.Moore, Z.Phys.C-Particle and Fields, **20** (1983), 71.

[25] B.W.Harris, J.Smith and R. Vogt, Nucl.Phys. **B 461**, (1996), 181.

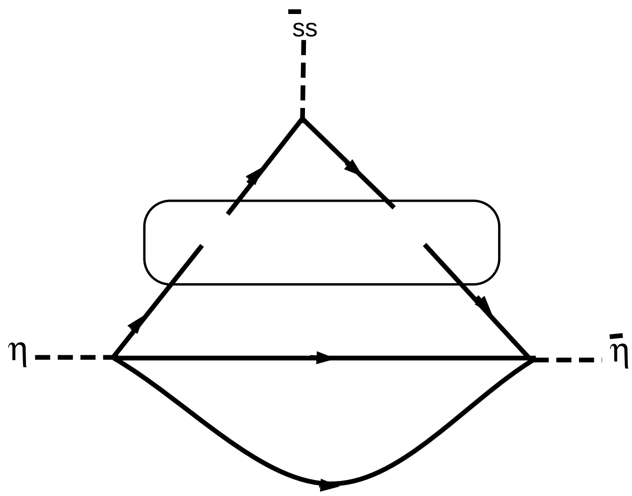

Fig.1

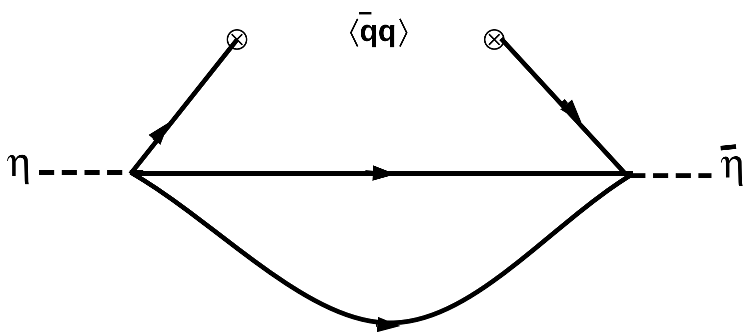

Fig.2